

GeV Emission from Collisional Magnetized Gamma Ray Bursts

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ABSTRACT

Magnetic fields may play a dominant role in gamma-ray bursts, and recent observations by *Fermi* indicate that GeV radiation, when detected, arrives delayed by seconds from the onset of the MeV component. Motivated by this, we discuss a magnetically dominated jet model where both magnetic dissipation and nuclear collisions are important. We show that, for parameters typical of the observed bursts, such a model involving a realistic jet structure can reproduce the general features of the MeV and a separate GeV radiation component, including the time delay between the two. The model also predicts a multi-GeV neutrino component.

Subject headings: gamma-rays: bursts — cosmology — stars: population III — jets: magnetized – radiation mechanisms: non-thermal

1. Introduction

Recent observations with the *Fermi* satellite have shown the presence of GeV emission in GRB bursts, delayed by seconds relative to the usual MeV radiation, e.g. Abdo et al. (2009a); Ackermann et al. (2010). In some bursts, this delayed GeV radiation appears as a distinct spectral component (Ackermann et al. 2010; Abdo et al. 2009b). There have also been indications that the outflow in such bursts may be magnetized, e.g. Zhang and Pe’er (2009); Fan (2009). Also on the theoretical side, there has been an increasing interest in magnetically dominated but baryon loaded GRB models, e.g. Thompson (1994); Drenkhahn and Spruit (2002); Lyutikov and Blandford (2003); Metzger et al. (2011); McKinney and Uzdensky (2011).

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The fast rotating central engine, whether a black hole or a magnetar, can lead to a highly magnetized outflow which is initially Poynting dominated, with a sub-dominant baryon load at the base of the flow r_0 such that $\eta \equiv (\mathcal{E}_{EM}/\rho_b c^2)_0 \gg 1$, where \mathcal{E}_{EM} is electromagnetic (mainly magnetic) energy density, and ρ_b is baryon mass density. If the magnetic field in the jet is striped, i.e. it has alternating polarity, either as a result of flow instabilities in a collapsar outflow or in oblique magnetar outflows, reconnection and dissipation can occur both below and above the photosphere, as discussed by Giannios and Spruit (2007); Metzger et al. (2011); McKinney and Uzdensky (2011), which can give rise to efficient radiation in the MeV band. Dissipative photospheres have also been proposed in magnetically sub-dominant baryon loaded outflows, as the source of the MeV radiation, e.g. Rees and Mészáros (2005); Pe’er et al. (2006); Ryde et al. (2010); Pe’er et al (2011); Beloborodov (2011), while the role of nuclear collisions in producing dissipation, high energy photons and neutrinos has been discussed by, e.g. Mészáros and Rees (2000a); Rossi et al. (2006); Ioka (2010); Beloborodov (2010) and others.

In this paper, we analyze a magnetically dominated and baryon loaded outflow, and discuss the combined effects of dissipation from magnetic reconnection and nuclear collisions in the context of the magnetically dominated outflow dynamics, which differs from the usual hydrodynamic regime dynamics. We show that such outflows can produce the usual MeV spectrum, as well as a GeV component which arrives with a delay of seconds after the MeV component, for a minimally (realistic) structured jet configuration.

2. Magnetic Dynamics, Nuclear Collisions and Photosphere

Dynamics.– If the magnetic field is striped and the jet is one-dimensional³ the outflow accelerates (Drenkhahn 2002; Granot et al. 2011; Metzger et al. 2011) roughly as

$$\Gamma(r) \simeq \begin{cases} (r/r_0)^{1/3} & \text{for } r \leq r_{sat} \\ \eta & \text{for } r \geq r_{sat} \end{cases} \quad (1)$$

where $r_0 = 10^7 r_{0.7}$ is the base of the outflow and $r_{sat} \simeq r_0 \eta^3$ is the radius at which the bulk Lorentz factor Γ has saturated, the outflow having become approximately matter dominated. One way to roughly understand this dynamic behavior is as follows. The comoving reconnection time is $t'_r \sim \lambda'/v'_r = \Gamma \lambda/v'_r \sim \Gamma \epsilon_1 r_0/\epsilon_2 v_A \propto \Gamma$, where $\lambda = \epsilon_1 r_0$ is the approximate radial lengthscale of field polarity changes, $\epsilon_1 \sim 1$, $v'_r \sim \epsilon_2 v'_A$ is reconnection velocity,

³For more general magnetic jet geometries and acceleration laws see text below eq.(13) and Komissarov et al. (2009); Tchekhovskoy et al. (2010); McKinney and Uzdensky (2011).

$v'_A \sim c$ is the Alfvén velocity, with, e.g., $\epsilon_2 \lesssim 0.1$ and comoving (lab) frame quantities are denoted as primed (unprimed). The dynamic expansion time is $t'_{ex} \sim r/c\Gamma$. The comoving internal energy per baryon of the plasma (which is mainly magnetic, in the acceleration phase) can be expressed as γ' . The outflow will be self-similar, and $\gamma' \geq 1$, a dimensionless quantity, can be expressed as a dimensionless combination of dimensional quantities in the flow. There are only two physically independent dynamical quantities, t'_r and t'_{ex} , suggesting $\gamma' \propto t'_r/t'_{ex} \propto \Gamma^2/r$. The conversion of internal energy during the expansion is what leads to the growth of the bulk kinetic energy, and energy conservation requires $\gamma'\Gamma \sim \text{const.} \propto \Gamma^3/r$, or $\Gamma \propto r^{1/3}$. The final bulk Lorentz factor is the same $\Gamma \simeq \eta$ as in the pure hydrodynamic regime, but the growth is slower, because at each e-folding in radius only part of the magnetic internal energy gets converted directly into kinetic energy, the other part gets converted first into thermal energy, and only after another e-folding into kinetic energy.

Longitudinal pion optical depth.– In a simple flow which initially moves with a single bulk Lorentz factor, nuclear inelastic collisions couple n and p with a cross section $\sigma_{nuc} \sim \sigma_\pi(c/v_{rel})$, where v_{rel} is the relative drift velocity between n and p and $\sigma_\pi \sim 3 \times 10^{-26} \text{ cm}^2$. When $v_{rel} \rightarrow c$, the nuclear collisions become inelastic resulting in the production of pions, muons, and eventually e^\pm pairs and neutrinos. The latter escape, but $\sim 1/4$ of the collision energy ends up in pairs and photons, which are trapped in the optically thick flow, adding to its thermal energy.

In such simple homogeneous flows, the nuclear collisions become inelastic (resulting in pions) when the comoving expansion time $t'_{ex} \sim r/c\Gamma$ becomes shorter than the comoving collision time $t'_{nuc} \sim 1/(n'_p \sigma_\pi c)$, or $\tau'_{\pi,||} \sim n'_p \sigma_{nuc} r / \Gamma \sim 1$. Here n'_p is the comoving proton density

$$n'_p = \frac{Lx}{4\pi r^2 m_p c^3 \eta \Gamma}, \quad (2)$$

where L is the total kinetic luminosity, and $x = n_p/n_b \equiv n_p/(n_p + n_n)$ is the proton fraction of the baryon density $n_b = n_p + n_n$. The radius at which pion production occurs is

$$\frac{r_\pi}{r_0} = \frac{L\sigma_\pi x}{4\pi m_p c^3 r_0 \eta \Gamma^2} \equiv \eta_\pi^6 \frac{x}{\eta \Gamma^2}, \quad (3)$$

where

$$\eta_\pi \equiv \left(\frac{L\sigma_\pi}{4\pi m_p c^3 r_0} \right)^{1/6} \simeq 1.33 \times 10^2 L_{54}^{1/6} r_{0,7}^{-1/6}, \quad (4)$$

and we assume parameters comparable to those deduced for typical bright LAT bursts, e.g. GRB 080916C (Abdo et al. 2009a), $L \sim 10^{54} L_{54}$, $\eta \gtrsim 600 \eta_{600}$, $\theta \sim 10^{-2}$.

The nuclear collisions can become inelastic in the regime $\Gamma \simeq (r/r_0)^{1/3} < \eta$, in which

case from eqs. (1, 3) the longitudinal pionization radius is

$$\frac{r_\pi}{r_0} = x^{3/5} \eta_\pi^3 \left(\frac{\eta_\pi}{\eta} \right)^{3/5} \quad \text{for } r_\pi < r_{sat}, \quad \text{i.e. } \eta > \eta_\pi x^{1/6} = 0.9 x_{0.5}^{1/6} \eta_\pi. \quad (5)$$

The pionization radius r_π and η_π of eqs. (5,4) scale with L , r_0 differently than in the corresponding hydrodynamic outflows, e.g. Bahcall and Mészáros (2000) and others, due to using here the magnetized outflow dynamics (1) instead of $\Gamma \propto r$. Beyond r_π , the neutrons decouple from the protons and start to coast with a terminal bulk Lorentz factor

$$\Gamma_n \simeq x^{1/5} \eta_\pi (\eta_\pi / \eta)^{1/5} \simeq 80 L_{54}^{1/5} \eta_{600}^{-1/5} x_{0.5}^{1/5}, \quad (6)$$

while the protons continue to accelerate as $\Gamma \propto r^{1/3}$ until reaching the coasting value $\Gamma \sim \eta$. If on the other hand the flow Lorentz factor saturates before the protons and neutrons decouple, i.e. before the nuclear collisions become inelastic, which occurs if $\eta < \eta_\pi x^{1/6}$, the protons and neutrons continue to coast together at $\Gamma \sim \eta = \text{constant}$ and pionization never occurs.

Photon optical depth.– If the only electrons in the outflow are those associated with the protons, the Thomson photosphere occurs at

$$\frac{r_{ph}}{r_0} = \left(\frac{L \sigma_T}{4\pi m_p c^3 r_0} \right) \frac{1}{\eta \Gamma^2} = \eta_T^6 \frac{1}{\eta \Gamma^2} \quad (7)$$

where

$$\eta_T = \left(\frac{L \sigma_T}{4\pi m_p c^3 r_0} \right)^{1/6} \equiv \eta_\pi \left(\frac{\sigma_T}{\sigma_\pi} \right)^{1/6} \simeq (5/3) \eta_\pi \simeq 2.22 \times 10^2 L_{54}^{1/6} r_{0,7}^{-1/6}. \quad (8)$$

For $\eta > \eta_T$ ($\eta < \eta_T$) the photosphere occurs at $r < r_{sat}$ ($r > r_{sat}$), at a radius

$$\frac{r_{ph}}{r_0} = \begin{cases} \eta_T^3 (\eta_T / \eta)^{3/5} & \text{for } r < r_{sat} \quad (\eta > \eta_T); \\ \eta_T^3 (\eta_T / \eta)^3 & \text{for } r > r_{sat} \quad (\eta < \eta_T). \end{cases} \quad (9)$$

The scaling, as well as the definition of η_T again differs from those in e.g. Mészáros and Rees (2000b) or Beloborodov (2010), due to our using here the magnetized outflow dynamics. Thus, the photosphere is typically larger then the longitudinal pionosphere by a factor

$$\frac{r_{ph}}{r_\pi} \sim \left(\frac{\sigma_T}{\sigma_\pi} \right)^{3/5} \sim 6.4. \quad (10)$$

The inelastic nuclear collisions produce additional leptons which can contribute to the Thompson opacity. For a strongly magnetized outflow the initial secondary leptons would cool very fast and the cascade efficiency drops as the pair production approaches its threshold,

where in the limit (Beloborodov 2010; Vurm et al. 2011) each n, p collision adds at most one further lepton. Since $\tau_\pi \propto r^{-5/3}$ drops rapidly, most neutrons will suffer only one inelastic collision, increasing the electron density by a factor ~ 2 . In addition, the reconnection process can also produce pairs, whose effects could be significant (McKinney and Uzdensky 2011; Metzger et al. 2011). Assuming that the two effects contribute a comparable change in the lepton density, this increases L by a factor $\gtrsim 4$ in eqs.(7,8, 9). Thus, the photospheric radius could be a factor $\gtrsim 4^{3/5} \sim 2.3$ larger than its value without pairs in eqs. (7,9), i.e. $r_{ph} \gtrsim 10 - 15r_\pi$.

Transverse pion optical depth.- Neutrons, if present outside the jet, can drift sideways into the jet channel unhindered by magnetic fields. The lab-frame transverse nuclear collision optical depth of a jet of opening half-angle θ at radius r is

$$\tau_{\pi,\perp} \sim n_p \sigma_\pi r \theta = \left(\frac{L \sigma_\pi}{4\pi m_p c^3 r_0} \right) \frac{r_0}{r} \frac{\theta}{\eta} \quad (11)$$

and the jet becomes transversely optically thin at

$$\frac{r_{\pi,\perp}}{r_0} = \eta_\pi^6 \frac{\theta}{\eta}. \quad (12)$$

Thus, for a jet opening half-angle $\theta = 10^{-2}\theta_{-2}$, we have the following radii,

$$\begin{aligned} r_\pi &\simeq 6 \times 10^{12} L_{54}^{3/5} r_{0,7}^{-3/5} x_{0.5}^{3/5} \eta_{600}^{-3/5} \text{ cm}, \\ r_{ph} &\simeq 6 \times 10^{13} L_{54}^{3/5} r_{0,7}^{-3/5} \eta_{600}^{-3/5} \text{ cm}, \\ r_{\pi,\perp} &\simeq 9 \times 10^{14} L_{54} \eta_{600}^{-1} \theta_{-2} \text{ cm}, \\ r_{sat} &\simeq 2 \times 10^{15} r_{0,7} \eta_{600}^3 \text{ cm} \end{aligned} \quad (13)$$

Transverse jet structure and neutron drift.- In a realistic jet the properties vary also in the transverse direction. Hydrodynamical simulations (e.g. (Morsony et al. 2010)) indicate a Lorentz factor tapering off towards the edges, while MHD force-free simulations (Tchekhovskoy et al. 2008) show models where, depending on the stellar pressure profile and the magnetic symmetries, the jet opening angle can initially narrow with increasing radius, with an inner jet core accelerating initially more slowly than the outer jet portions, but before exiting the star the inner jet becomes faster than the outer sheath, after which the jet opening angle remains constant. As a simple idealized model, we can consider the transverse structure outside the star as a two-step jet, consisting of an inner jet core with, e.g., a nominal constant half-angle $\theta = 10^{-2}\theta_{-2}$ with $\eta = 600\eta_{600}$, and a slower⁴ outer jet or

⁴Depending on various parameters an alternative model might have outside the star a slower inner jet and a faster sheath, which could lead to results similar to those discussed below; we do not consider this case here.

sheath with, say, $\eta_{out} \sim 100$, where both inner core and outer sheath have been populated with protons and free neutrons already near the black hole, where most nuclei get photo-dissociated (Beloborodov 2003; Metzger et al. 2008). One can expect neutrons from the outer sheath to drift into the jet core, and for $\eta_{out} \lesssim 10^2 \sim \theta^{-1}$ the neutrons can penetrate transversely the entire jet core. Also, the relative radial Lorentz factor ratio between indrift neutrons and jet core baryons (either p or n) is $\Gamma_{rel} \sim \eta/\eta_{out} > 1$, so the collisions will be inelastic, leading again to pions.

Over a timescale t the thermal neutrons from an outer sheath where the lab-frame neutron density is n_n and the thermal neutron velocity is v_t will diffuse transversely into the jet, their flux being

$$\phi_n \sim \left(\frac{v_t \ell}{t}\right)^{1/2} n_n \sim \beta_t \left(\frac{n_n c}{\sigma_\pi t}\right)^{1/2}, \quad (14)$$

where $\ell \sim 1/(n_n \sigma_{nuc})$, $\sigma_{nuc} \sim \sigma_\pi (c/v_t)$ and $\beta_t = v_t/c$. Taking a jet transverse area of $A \sim \pi \theta r^2$, on a timescale t the total number of neutrons drifting transversely into the jet core will be, roughly,

$$N_{n,\perp} \sim \pi \theta r^2 \phi_n t \sim \pi \theta r^2 \beta_t (n_n c t / \sigma_\pi)^{1/2}, \quad (15)$$

while over the same timescale the number of baryons (n and p) passing longitudinally through the jet core is

$$N_{b,\parallel} \sim \pi \theta r^2 n_b c t \sim \pi \theta r^2 (L/4\pi r^2 m_p c^3 \eta) c t \quad (16)$$

The number of collisions is maximized when $N_{n,\perp} \gtrsim N_{b,\parallel}$, that is $\beta_t (n_n c t / \sigma_\pi)^{1/2} \gtrsim (c t / \sigma_\pi) (\eta_\pi^6 r_0 / r^2 \eta)$, and when the jet transverse pion optical depth $\tau_{\pi,\perp} \sim 1$, i.e. $r \sim r_{\pi,\perp}$. For the first condition, we can estimate the outer sheath neutron thermal velocity β_t since we know that the comoving baryon temperature becomes trans-relativistic ($T \sim \text{GeV}$) at the saturation radius $r_{sat,out} = r_{0,7} \eta_{out}^3 = 10^{13} \text{ cm}$, where $v_t \sim c$, after which they coast with $T_t \propto r^{-2/3}$, so at $r_{\pi,\perp}$ we have $\beta_t^2 = (v_t/c)^2 \sim (k T_t / m_p c^2) = (r_{sat,out} / r_{\pi,\perp})^{-2/3} \sim 10^{-4/3}$. These two conditions are then met when at $r \sim r_{\pi,\perp}$ the neutron density in the outer sheath is

$$n_n \gtrsim \frac{c t}{\sigma_\pi} \frac{\eta^2}{\eta_\pi^{12}} \frac{1}{r_0^2 \theta^4 \beta_t^2} \sim 4 \times 10^{11} r_{0,7}^{-2} \theta_{-2}^4 L_{54}^{-2} \eta_{600}^2 t \text{ cm}^{-3} \quad (17)$$

over a timescale t . The density in eq. (17) is comparable to what might be expected in an outer jet sheath.

3. Spectrum Formation

GeV radiation.- The ‘prompt’ component of the GeV radiation is expected to arise from the transverse drift nuclear collision mechanism at $r_{\pi,\perp}$ discussed above. The $\gamma\gamma$ photo-

sphere to $E = 10$ GeV photons can be roughly estimated as (Beloborodov 2010) $r_{\gamma\gamma}(E) \sim 10^{15}(E/10 \text{ GeV})^{3/2}\eta_{600}^{-5}L_{54} \text{ cm}$, so that

$$r_{ph} < r_{\gamma\gamma}(10 \text{ GeV}) \lesssim r_{\pi,\perp}, \quad (18)$$

the right inequality holding for $\eta_{600} \gtrsim 1$. Thus, at $r_{\pi,\perp}$, which is at or outside $r_{\gamma\gamma}$ for $\eta_{600} \gtrsim 1$, multi-GeV photons will be copiously produced by transverse indrift neutron collisions with jet core baryons, and these photons can escape unhindered. The energy in this GeV component, for the conditions discussed in the previous section, is a significant fraction of the kinetic luminosity. Most pions will be formed approximately at rest in the frame of the jet, and the processes $p, N \rightarrow \pi^+, \pi^-, \pi^0$, where $N = (p, n)$, occur in approximately equal ratios. Thus 1/3 of the collisions on average yield two photons from $\pi^0 \rightarrow 2\gamma$ decays at an observer frame energy which is broadly distributed around a central energy $\varepsilon_{\gamma}^{\pi,\perp} \sim \eta(m_{\pi}c^2/2)/(1+z) \sim 7 \eta_{600}(5/1+z) \text{ GeV}$, extending up to twice that value and down to a fraction of it.

GeV time delay.- The relative time delay between between MeV photons from the central photosphere (see below) and the observed onset of the GeV emission, which arises at $r_{\pi,\perp}$ from neutrons from an outer jet sheath moving with bulk Lorentz factor η_{out} which drift into the inner jet at the radius $r_{\pi,\perp}$ where their effect becomes significant is

$$\Delta t \sim \frac{r_{\pi,\perp}}{2c\eta_{out}^2} \sim 1 L_{54}r_{0,7}\eta_{600}^{-1}\eta_{out,100}^{-2} \text{ s}, \quad (19)$$

in the lab frame at the source. In the observer frame this is lengthened by a factor $(1+z)$. For a burst at $z \sim 4$ such as GRB 080916C this is similar to the observed GeV-MeV delay of ~ 5 seconds observed by Fermi in that object.

MeV radiation.- The MeV (Band) spectrum may be expected to arise from the dissipative photosphere of the inner jet at $r_{ph} \sim 6 \times 10^{13} \text{ cm}$ (eq.[13]). This MeV spectrum may be attributed to synchrotron radiation following magnetic or shock dissipation near r_{ph} , with typical bulk Lorentz factors $\Gamma_r \sim 1$. This would result in random magnetic fields $B' \sim (32\pi\epsilon_B m_p c^2 n'_b)^{1/2}\Gamma_r$, where $n'_b = L/(4\pi r_{ph}^2 \eta \Gamma_{ph})$ is the baryon comoving density at r_{ph} , where for the nominal jet parameters used we have $\Gamma_{ph} = (r_{ph}/r_0)^{1/3} \sim 180 L_{54}^{1/5} r_{0,7}^{-8/15} \eta_{600}^{-1/5}$, so $n'_b \sim 5 \times 10^{12} L_{54}^{-2/5} r_{0,7}^{26/15} \eta_{600}^{2/5}$ and $B' \sim 4 \times 10^5 L_{54}^{-1/5} r_{0,7}^{13/15} \eta_{600}^{1/5} \epsilon_B^{1/2} \Gamma_r \text{ G}$. The same magnetic dissipation leading to semi-relativistic shocks with comoving bulk Lorentz factors $\Gamma_r \sim 1$ near r_{ph} will accelerate electrons to average non-thermal random Lorentz factors of $\gamma'_e \sim (m_p/m_e)\Gamma_r \sim 10^3 \Gamma_r$, assuming equipartition between baryons and electrons, leading to an observer frame synchrotron radiation peak at $\varepsilon_{sy}^{ph} = \varepsilon'_{sy} \Gamma_{ph} (1+z)^{-1} \sim 0.3 r_{0,7}^{1/3} \epsilon_B^{1/2} \Gamma_r^3 (5/1+z) \text{ MeV}$. This non-thermal component will dominate over a much weaker photospheric thermal component at $kT_{th} \sim 1 \text{ keV}$. Comptonization may boost a fraction of the synchrotron photons to higher energies, but using the expression above for $r_{\gamma\gamma}$ a cutoff is obtained at

$\varepsilon_{\gamma\gamma} \sim 3 \text{ MeV} L_{54}^{-2/3} \Gamma_{ph,180}^{10/3}$. Thus, the photosphere allows escape of MeV but not GeV photons. A decreasing fraction of increasing energy photons could arise from dissipation at larger radii (Beloborodov 2010), as $r_{\gamma\gamma}$ increases, but significant GeV escape compatible with a second spectral component is only expected from radii where transverse neutron drift is important.

Outer jet photosphere.– The outer jet has a photosphere beyond its saturation radius, at $r_{ph,out} \sim 2.5 \times 10^{14} L_{54} \eta_{out,100}^{-3}$ cm. If there was dissipation at this outer photosphere, a calculation similar to the inner jet’s would give a synchrotron peak at $\varepsilon_{sy,out} \sim 25 L_{54}^{-1/2} \epsilon_B^{1/2} \Gamma_r^3 \eta_{out,100}^{1/2} (5/1+z)$ keV. However, magnetic dissipation is not expected to occur above the saturation radius, and neither does nuclear collisional dissipation, since the n, p components do not decouple for this $\eta < \eta_\pi \sim 130$. Thus, unless some other form of dissipation occurs near the outer photosphere, it would have a quasi-blackbody spectrum with $kT \sim 0.2 L_{54}^{-1/2} r_{0,7}^{1/6} \eta_{out,100}^{2/3} (5/1+z)$ keV and a luminosity $L_{ph,out} \sim \epsilon_{rad} L_0 (r_{sat}/r_{ph})^{2/3} \lesssim 10^{-2} (\epsilon_{rad}/0.5) L_0$, well below the kinetic luminosity L_0 .

Internal shocks.– In principle these would not be expected if the magnetization remains significant beyond the saturation radius (c.f. Zhang and Yan (2011)). If they did occur, the earliest would need to be at $r_{IS} \sim 2ct_v \eta^2 \gtrsim 2 \times 10^{15} t_{v,-1} \eta_{600}^2$ cm $\gtrsim r_{sat}$, where $t_v = 10^{-1} t_{v,-1}$ s is the variability time, and the usual arguments would lead to a synchrotron peak at $\varepsilon_{sy}^{IS} \lesssim 8 L_{54}^{1/2} \epsilon_B^{1/2} \epsilon_e^2 t_{v,-1}^{-1} \eta_{600}^{-1} \Gamma_r (5/1+z)$ keV.

External shocks and afterglow.– Other radiation components are due to the external shocks of the inner and outer jets. Forward shocks (FS) are guaranteed, while reverse shocks (RS) are less certain, depending on the residual magnetization σ value at the external shock. For an outflow (prompt burst) duration $t = 10t_1$ s in the burst frame, in an external medium of density $n_0 \text{ cm}^{-3}$, the inner jet external shock occurs at a deceleration radius $r_d \sim 1.6 \times 10^{17} L_{54}^{1/3} t_1^{1/3} n_0^{-1/3} \eta_{600}^{-2/3}$ cm at an observer time $t_d \sim 50 L_{54}^{1/3} t_1^{1/3} n_0^{-1/3} \eta_{600}^{-8/3} ([1+z]/5)$ s. With the usual assumption about magnetic field and electron energy equipartition, and $\Gamma_d \sim \eta/2$ at r_d , the FS synchrotron spectrum of the inner jet peaks at $\varepsilon_{sy}^{FS} \sim 2\epsilon_B^{1/2} \epsilon_e^2 \eta_{600}^4 (5/1+z)$ MeV. For the usual power law electron distribution accelerated at the shock, this synchrotron spectrum extends as a power law into the GeV range (Ghisellini et al. 2010; Kumar and Barniol 2010), providing the long term GeV afterglow, the external shock being well outside $r_{\gamma\gamma}(10 \text{ GeV})$. The earliest prompt GeV radiation, however, appears to require a separate origin, as considered by Toma et al. (2011); He et al. (2011); De Pasquale et al. (2010); Corsi et al. (2010) in the context of non-magnetic outflows.

The outer jet deceleration radius is $r_{d,out} \sim 5.3 \times 10^{17} L_{54}^{1/3} t_1^{1/3} n_0^{-1/3} \eta_{100,out}^{2/3}$ cm at an observer time $t_{d,out} \sim 2 \times 10^3 L_{54}^{1/3} t_1^{1/3} n_0^{-1/3} \eta_{100,out}^{-8/3} ([1+z]/5)$ s, and the outer jet FS synchrotron

spectral peak is at $\varepsilon_{sy,out}^{FS} \sim 2 \epsilon_B^{1/2} \epsilon_e^2 \eta_{out,100}^4 (5/1 + z)$ keV.

If the magnetization parameter of the ejecta σ has become small at r_d , a reverse shock may form (Mimica et al. 2009; Zhang and Yan 2011). With the usual assumptions for an equipartition magnetic field in a reverse shock with $\Gamma_r \sim 1$ and $\gamma'_e \sim \epsilon_e 10^3 \Gamma_r$ the inner jet RS synchrotron peak occurs at $\varepsilon_{sy}^{RS} \sim 70 L_{54}^{1/6} t_1^{1/3} n_0^{1/2} \epsilon_B^{1/2} \epsilon_e^2 \Gamma_r^3 \eta_{600}^{2/3} (5/1 + z)$ eV, while the outer jet reverse shock synchrotron spectrum peaks at $\varepsilon_{sy,out}^{RS} \sim 20$ eV, with the same scaling.

Neutrino spectrum.– Also, as a result of the charged pion decays, 1/3 of the collisions yield on average a muon neutrino of energy broadly centered around $\varepsilon_\nu \sim 7 \eta_{600} (5/1 + z)$ GeV (again extending up to twice that value and down to a fraction of it), and a muon neutrino as well as an electron neutrino centered at energy $\varepsilon_\nu \sim 3 \eta_{600} (5/1 + z)$ GeV, while another 1/3 of the collisions leads to the corresponding antineutrinos of similar energies. These are in the sensitivity range of the Deep Core array of the IceCube neutrino detector.

4. Discussion

We have presented a GRB model based on a magnetically dominated, baryon loaded outflow, where both magnetic reconnection and nuclear p, n collisions lead to dissipation around and above a photon photosphere. We have discussed the dynamics of the magnetized jets based on a simplified reconnection prescription and an idealized jet transverse structure, deriving the radii at which the photon scattering photosphere occurs as well as the radii at which radial and transverse drifts between proton and neutron components lead to nuclear collisional dissipation, in the context of the magnetic jet dynamics. Although the detailed structure and dynamics could depend on various assumptions about the stellar envelope pressure profile, the magnetic field configuration, etc. (Tchekhovskoy et al. 2008; McKinney and Uzdensky 2011; Metzger et al. 2011), the present model qualitatively captures a number of the essential features. In particular, the characteristic radii and their scalings differ from those in a purely hydrodynamical jet. Both in this and in more general magnetic models a slower acceleration rate generically leads to larger photospheric and dissipation radii relative to the hydrodynamic case, in line with previous hints from *Fermi* observational results indicating larger emission radii.

Using parameters inferred from *Fermi* observations of bursts showing significant emission in the LAT instrument in the GeV energy range, which typically are of very high luminosity and higher than average Lorentz factor (Zhang et al. 2011), we discussed the gross properties of the MeV and GeV photon production spectra. For a simplified structured jet model consisting of a central faster jet core and a slower outer jet surrounding it, we showed that for

reasonable parameters a significant GeV component is produced, with a time delay of seconds relative to the onset of the MeV component, in rough agreement with the observations of GRB 080916C, e.g. (Abdo et al. 2009a). Related to the photon emission, a multi-GeV neutrino emission component is also expected, whose luminosity is comparable to that of the multi-GeV photon emission. More detailed calculations, which are beyond the object of this letter, will be needed in order to make detailed spectral fits to specific bursts.

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